Is the Higgs Mechanism of Fermion Mass Generation a Fact? A Yukawa-less First-Two-Generation Model

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It is now established that the major source of electroweak symmetry breaking (EWSB) is due to the observed Higgs particle. However, whether the Higgs mechanism is responsible for the generation of all the fermion masses, in particular, the fermions of the first two generations, is an open question. In this letter we present a construction where the light fermion masses are generated through a secondary, subdominant and sequestered source of EWSB. This fits well with the approximate U(2) global symmetry of the observed structure of the flavor sector. We first realise the above idea using a calculable two Higgs doublet model. We then show that the first two generation masses could come from technicolor dynamics, while the third generation fermions, as well as the electroweak gauge bosons get their masses dominantly from the Higgs mechanism. We also discuss how the small CKM mixing between the first two generations and the third generation, and soft mixing between the sequestered EWSB components arise in this setup. A typical prediction of this scenario is a significant reduction of the couplings of the observed Higgs boson to the first two generation of fermions.

The mechanism of Electroweak Symmetry Breaking (EWSB) and the generation of fermion masses are among the main frontiers of research in high energy physics. The discovery of a Higgs-like particle and the measurements of its couplings to the Standard Model (SM) gauge bosons have established that the Higgs mechanism is indeed the correct picture for it, at least to leading order. However, whether the Higgs particle is also responsible for all the fermion masses is an open question. We have direct indications that the observed Higgs particle couples to the third generation fermions with roughly SM strengths

$$\mu_{t\bar{t}h,\,b,\,\tau} = 2.2 \pm 0.6,\ 0.71 \pm 0.31,\ 0.97 \pm 0.23\,, \quad (1)$$

where we have averaged the ATLAS [1–3] and CMS [4–6] results for the corresponding signal strengths. The couplings to the first two generation of fermions (also referred as light fermions below) have not been measured and are only weakly constrained: [7–10]

$$\mu_\mu \le 7\,,\ \mu_e \le 4\times 10^5\,,\ \mu_c \le 180\,, \quad (2)$$

while for the u, d, s quarks no direct bounds exist at present (see [11–17] for related discussions). In addition, the observed flavor sector consists of a large hierarchy between the first two and the third generation fermion masses and the mixing angles. This is related to the celebrated SM flavor puzzle that is consistent with an approximate pattern of U(3)/U(2) symmetry breaking [18, 19].

Below we explore the possibility that the above approximate U(2) structure is linked with the way EWSB is communicated to the flavor sector. We propose that two approximately sequestered sources of EWSB exist in nature [20]. The major one, is due to a SM-like Higgs doublet, Φ_3, that couples predominantly to the third generation fermions. An additional subdominant source, Φ_{12}, couples mostly to the first two generations and induces their masses. The second source of EWSB Φ_{12} has very little to do with the observed Higgs particle. It could be a cousin of the Higgs or could arise due to strong-dynamics, and not be associated with any weakly coupled physics. More concretely, we assume that Φ_3 has a mass of about 125 GeV with the following couplings

$$\mathcal{L}_3^Y \supset -\overline{Q}_L^{(i)}[Y^d_{3\ ij}\,\Phi_3\, d_R^{(j)} + Y^u_{3\ ij}\,\widetilde{\Phi}_3\, u_R^{(j)}] + \text{h.c.}$$

with $Y^{u,d}_{3\ ij} \approx y^{t,b}\delta_{33}$. Assuming that the other source of EWSB transforms as a doublet (to preserve custodial symmetry), Φ_{12}, the mass terms of the light fermions can be effectively written as,

$$\mathcal{L}_{12}^Y \supset -\overline{Q}_L^{(i)}[Y^d_{12\ ij}\,\Phi_{12}\, d_R^{(j)} + Y^u_{12\ ij}\,\widetilde{\Phi}_{12}\, u_R^{(j)}] + \text{h.c.}$$

where i, j run from 1 to 2 and,

$$v_{12}^2 = \langle\Phi_{12}\rangle^2 \ll v_3^2 = \langle\Phi_3\rangle^2\,,\quad v_{12}^2 + v_3^2 = v^2\,, \quad (3)$$

with $v = 246$ GeV. Later, we will explore the possibility that Φ_{12} is actually a condensate, $\langle\bar{Q}Q\rangle$, of fermions of a technicolor sector. If the sector containing Φ_3 and the third generation of fermions is completely decoupled from the second source of EWSB,

the observed Higgs boson would have no couplings to the light fermions.

Two conceptual problems arise, however, if the two EWSB sectors are completely sequestered. First, the (13) and (23) elements of the Cabibbo-Kobayashi-Maskawa (CKM) matrix are not generated. In order to generate these CKM entries, we will assume that Φ_3 couples very weakly to the light fermions,

$$Y^{u,d}_{3\,ij} = y^{t,b}\delta_{33} + \epsilon^{u,d}_{ij}\,, \tag{4}$$

where, as shown below, the strength of the ϵ_{ij} couplings are dictated by the observed CKM (13) and (23) elements. The second issue is that, if there is no coupling between the two sectors at all, we can define two different global SU(2) symmetries for Φ_{12} and Φ_3. As a result, when Φ_3 and Φ_{12} develop Vacuum Expectation Values (VEVs), both the SU(2) symmetries are broken and six Goldstone bosons arise in the low energy spectrum. Three of these Goldstone bosons are absorbed by $W^\pm_L$ and Z_L, once the vectorial combination of the two SU(2) global symmetries is gauged. As the gauging explicitly breaks the axial combination of the two SU(2) global symmetries, the remaining three Goldstone bosons get small masses at the loop level. However, these masses are too small to evade collider constraints and a further source of explicit breaking is required to raise their masses. For this, we add, in our scalar potential, a μ-term,

$$\mu\,\Phi_3^\dagger\,\Phi_{12} + \text{h.c.}\,, \tag{5}$$

that breaks this symmetry softly, and lifts the Goldstone boson masses. Note that at the spurionic level, once we have the couplings ϵ_{ij} in Eq. (4) the coupling, μ, is also allowed. This means that generically, a mechanism that generates the ϵ_{ij} would also induce the μ-term. In our technicolor construction below, this is indeed the case.

With the addition of the term in Eq. (4) the light fermions get masses from two EWSB sources, thus violating a "Natural Flavor Conservation" principle [21]. This is expected to lead to tree level Flavor Changing Neutral Current (FCNC) processes, suppressed by powers of ϵ_{ij} (see for eg., [22]). We will see that, in order to generate CKM elements with the correct magnitude, the size of ϵ_{ij} needed is small enough such that FCNCs are within experimental bounds. In this sense our model is similar in spirit to other models that violate NFC Ref. [23–29] but nevertheless satisfies flavor bounds.

In order to understand these issues more concretely, we now consider an effective 2HDM toy model for Φ_3 and Φ_{12}, which will allow us to understand the relevant phenomenological constraints without committing to a specific UV completion. The following scalar potential is consistent with the setup described above [30],

$$\begin{aligned} V(\Phi_{12},\Phi_3) = \mu_1^2|\Phi_{12}|^2 + \mu_2^2|\Phi_3|^2 + \mu(\Phi_{12}^\dagger\Phi_3 + \text{h.c.}) \\ +\lambda_1|\Phi_{12}|^4 + \lambda_2|\Phi_3|^4\,. \end{aligned} \tag{6}$$

The couplings μ_1, μ_2, λ_1, λ_2 are real due to hermiticity of the Lagrangian. Moreover, the phase of μ can be absorbed in Φ_3. We have in the spectrum a pseudoscalar Higgs A and two charged Higgs bosons $H^\pm$ which model the pseudo goldstone bosons discussed above Eq. (5). There are also two CP-even scalars h and H. The masses and the mixing angle, β, for the pseudoscalar and charged states are given by, $m^2_{H^\pm} = m^2_{A^0} = -\mu v^2/v_{12}v_3$, $\tan\beta = v_3/v_{12}$ while the masses and the mixing angle, α, for the CP-even states are given by,

$$\begin{aligned} m^2_{h,\,H} &= \lambda_1 v_{12}^2 + \lambda_2 v_3^2 - \frac{\mu v^2}{2\,v_{12}\,v_3} \\ &\quad\mp \sqrt{\left(\lambda_1 v_{12}^2 - \lambda_2 v_3^2 + \frac{\mu(v_{12}^2 - v_3^2)}{2\,v_{12}\,v_3}\right)^2 + \mu^2}\,, \\ \tan 2\alpha &= -\frac{2m^2_{A^0}s_\beta c_\beta}{\sqrt{(m_H^2 - m_h^2)^2 - 4m^4_{A^0}s^2_\beta c^2_\beta}}\,, \end{aligned} \tag{7}$$

where $c_\theta \equiv \cos\theta$ and $s_\theta \equiv \sin\theta$. The coupling of the quarks with the Higgs bosons can be written as,

$$\begin{aligned} \mathcal{L}_Y^{(q)} \supset &-\frac{h}{v}\Big\{\bar u\Big[(s_{\beta-\alpha}M_u + c_{\beta-\alpha}Y_u)P_R + \\ &\qquad (s_{\beta-\alpha}M_u + c_{\beta-\alpha}Y_u^\dagger)P_L\Big]u + u \leftrightarrow d\Big\} \\ &- \Big(h \to H, s_{\beta-\alpha} \to c_{\beta-\alpha}, c_{\beta-\alpha} \to -s_{\beta-\alpha}\Big) \\ &+ \frac{iA}{v}\Big[\bar u(Y_u P_R - Y_u^\dagger P_L)u - \bar d(Y_d P_R - Y_d^\dagger P_L)d\Big] \\ &+ \Big[\frac{H^+}{v}\bar u(Y_u^\dagger V_{\text{CKM}}P_L - V_{\text{CKM}}Y_d P_R)d + \text{h.c.}\Big], \end{aligned} \tag{8}$$

where the matrices Y_u and Y_d are defined as,

$$Y_{u,d} = V_L^{u,d\,\dagger}\big(-v_3 Y_{12}^{u,d} + v_{12}Y_3^{u,d}\big)V_R^{u,d}/\sqrt{2}\,, \tag{9}$$

and the matrices V_L^u, V_L^d, V_R^u and V_R^d are defined through the following equations,

$$M_{u,d} = V_L^{u,d\,\dagger}(v_{12}Y_{12}^{u,d} + v_3 Y_3^{u,d})V_R^{u,d}/\sqrt{2}\,. \tag{10}$$

Here, M_u and M_d are the (diagonal) mass matrices. As for $Y_{12}^{u,d}$ and $Y_3^{u,d}$, without loss of generality we can assume that Y_{12}^d has the form,

$$Y_{12}^d = \text{diag}(y_d, y_s, 0) \tag{11}$$

where $y_{d,s} = \sqrt{2} m_{d,s}/v_{12}$ upto $\mathcal{O}(\epsilon^2)$ corrections. In the same basis,

$$Y_3^d = \begin{pmatrix} \mathcal{O}(\epsilon_{L1}^{d*}\epsilon_{R1}^d) & \mathcal{O}(\epsilon_{L1}^{d*}\epsilon_{R2}^d) & \epsilon_{L1}^{d*} \\ \mathcal{O}(\epsilon_{L2}^{d*}\epsilon_{R1}^d) & \mathcal{O}(\epsilon_{L2}^{d*}\epsilon_{R2}^d) & \epsilon_{L2}^{d*} \\ \epsilon_{R1}^d & \epsilon_{R2}^d & y_b \end{pmatrix} . \tag{12}$$

where $y_b = \sqrt{2} m_b/v_3$. In the above basis, Y_{12}^u is misaligned from the up sector mass basis by an angle $\mathcal{O}(V_{\text{CKM}}^{12})$. In the basis where $Y_{12}^u = \text{diag}(y_u, y_c, 0)$ is diagonal we define Y_3^u as in Eq. (12) but with the replacement $d \to u$ and $y_b \to y_t$. We take $\epsilon_{Li}^u \sim \epsilon_{Li}^d$.

Assuming the above form, we diagonalise $M_{u,d}$ (see Eq. (10)) and obtain $V_L^{u,d}$ and $V_R^{u,d}$. Then using $V_{\text{CKM}} = V_L^{u\dagger} V_L^d$ we find that the $\epsilon_{Li}^{u,d}$ are related to the CKM matrix elements as follows,

$$\begin{aligned} \epsilon_{L2}^u \sim \epsilon_{L2}^d &= \hat{y}_b V_{\text{CKM}}^{23}/s_\beta \, , \\ \epsilon_{L1}^u \sim \epsilon_{L1}^d &= \hat{y}_b V_{\text{CKM}}^{13}/s_\beta - \hat{y}_b V_{\text{CKM}}^{12} V_{\text{CKM}}^{23}/s_\beta \end{aligned} \tag{13}$$

where $\hat{y}_f = \sqrt{2} m_f/v$. Note that we still have freedom to choose $\epsilon_R^{u,d}$. Using Eq. (9) we find that the matrix Y_q, which governs the tree level FCNC interactions of the pseudoscalar and the charged Higgs bosons, to be

$$\begin{aligned} Y_q &= \cot\beta\, M_q - \\ &\frac{v}{\sqrt{2} c_\beta} \begin{pmatrix} \mathcal{O}(\epsilon^4) & \epsilon_{L1}^q \epsilon_{R2}^{q*} s_\beta/\hat{y}_3 & s_\beta^2 \epsilon_{L1}^q \epsilon_{R2}^{q*\,2}/\hat{y}_3^2 \\ \epsilon_{L2}^{q*}\epsilon_{R1}^{q*} s_\beta/\hat{y}_3 & \hat{y}_2/s_\beta & \epsilon_{R2}^{q*}\hat{y}_2/\hat{y}_3 \\ s_\beta^2 \epsilon_{R1}^{q*}\epsilon_{L2}^{q\,2}/\hat{y}_3^2 & \epsilon_{L2}^q \hat{y}_2/\hat{y}_3 & s_\beta \hat{y}_2 \epsilon_{L2}^q \epsilon_{R2}^{q*}/\hat{y}_3^2 \end{pmatrix} \end{aligned} \tag{14}$$

where $q = u, d$, $\hat{y}_3 = \hat{y}_{t,b}$ and $\hat{y}_2 = \hat{y}_{c,s}$. As expected Y_q has ϵ_{ij} suppressed off-diagonal entries which lead to Higgs mediated FCNCs. For example, the $\Delta F = 2$ operators $Q_2 = \bar{d}_R^\alpha q_L^\alpha \bar{d}_R^\beta q_L^\beta$, $\tilde{Q}_2 = \bar{d}_L^\alpha q_R^\alpha \bar{d}_L^\beta q_R^\beta$, $Q_4 = \bar{d}_R^\alpha b_L^\alpha \bar{d}_L^\beta b_R^\beta$ are generated by tree level exchange of the h , H and A. While the constraints from $B-\bar{B}$ mixing are trivially satisfied (due to the approximate alignment between Y_d and M_d which is only broken by y_s effects), contribution to ΔM_K is just below the experimental sensitivity for $\epsilon_{Li} \sim \epsilon_{Ri}$. Reducing the ratio $r_{RL}^\epsilon \equiv \epsilon_{Ri}/\epsilon_{Li}$ one can further suppress the contribution to Kaon mixing. If $\mathcal{O}(1)$ phases are present then the bound from ϵ_K requires $r_{RL}^\epsilon \lesssim 0.1$.

Note that we could have also generated the (13) and (23) elements of the CKM matrix by adding a perturbation, ϵ_{ij} to $Y_{12}^{u,d}$, instead of $Y_3^{u,d}$. In this case, however, it turns out that contributions to $B-\bar{B}$ mixing are a bit larger than the allowed values. This is because here, unlike the previous case, the approximate alignment between Y_d and M_d for $y_s \to 0$ is not present.

Let us now discuss constraints from Electroweak Precision Observables (EWPO) and other phenomenological implications. Following the expressions given in [27], we compute the Peskin-Takeuchi parameters S, T and U [31, 32] and show the allowed region in Fig. 1. The constraints from EWPO give an upper bound on v_{12}. For example, taking $m_H = 450$ and $m_A = 250$ GeV we find the bound $v_{12} < 80$ GeV. To evade the direct constraints from the Drell-Yan production of $H^\pm$ at LEP one also requires $m_{H^+} \gtrsim 100$ GeV, which, in turn, provides a lower bound of $\mu \gtrsim 50-80$ GeV for $v_{12} \approx 80-160$ GeV. Note that, while the couplings of the observed Higgs boson to the electroweak gauge bosons see a reduction of $\mathcal{O}(v_{12}^2/v^2)$ with respect to the SM value, couplings to the third generation fermions get an enhancement of the same order. Thus we need precision Higgs measurements to reach a $\mathcal{O}(10\%)$ sensitivity to detect these deviations.

Using Eq. (8) we also find that κ_f, the coupling of the observed Higgs boson to a light fermion f normalised to its SM value m_f/v, can be written as,

$$\kappa_f = -\frac{s_\alpha}{c_\beta} + \mathcal{O}(\epsilon^2) \xrightarrow{m_A \ll m_H} s_\beta \frac{m_A^2}{m_H^2} . \tag{15}$$

Hence, for $m_H^2 \gg m_A^2$ these couplings can be considerably reduced with respect to the SM. In order to quantify the extent of reduction in the Higgs coupling to the light fermions in our 2HDM framework, in Fig. 1 we show (in blue) the allowed region in the $m_A - m_H$ plane (for fixed v_{12}) and the $v_{12} - m_H$ plane (for fixed m_A) overlayed with contours of κ_f. One can notice that a factor of $2-3$ reduction in the coupling is possible. This is, however, at the cost of large values of one of the quartic couplings, λ_1 (λ_2 is almost fixed by the observed Higgs boson mass).

Thus, in the region of the parameter space, where the observed Higgs couplings to the light fermions are reduced, a large value of λ_1 renders the theory nonperturbative at the TeV scale. In this limit, our setup finds a natural embedding in a strongly coupled theory like technicolor where the VEV of Φ_{12} is identified with $\sqrt{N_D} F_T$, F_T and N_D being the technipion decay constant and the number of TC doublets respectively. The scalars A^0 and $H^\pm$ are identified with the (pseudo) goldstone bosons of the TC chiral

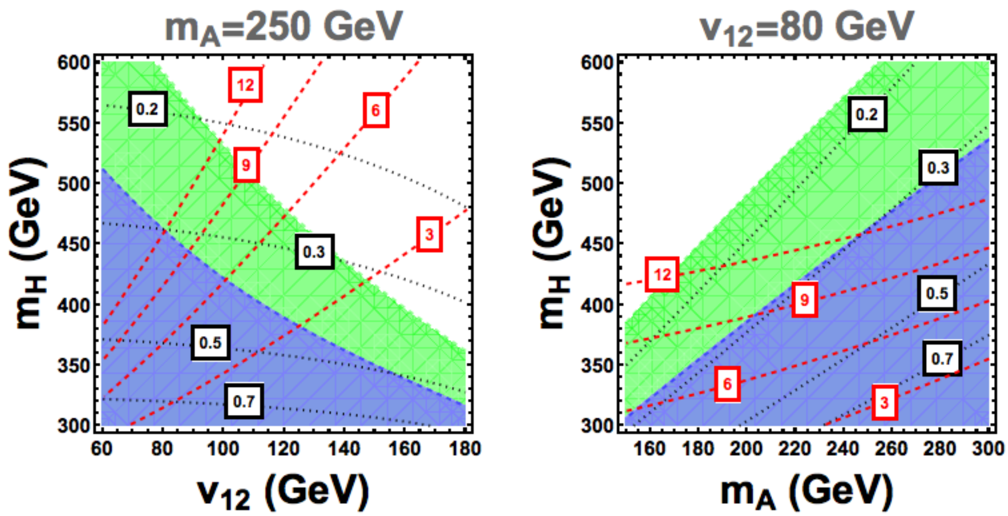

FIG. 1: In blue we show the region allowed by Electroweak precision constraints at 95% CL considering only the 2HDM contribution. The green region shows the improvement when an additional UV contribution {S,T} ={0.25,0.25} is added. The red and black dashed lines are contours of fixed λ_1 and fixed κ_f respectively. Note that, as the signal strength $\mu_f = \kappa_f^2$, the reduction in the signal strength is more significant.

symmetry breaking. The mass of the heavy CP-even Higgs, m_H, should be identified with the scale of the resonances,

$$m_{\rho_T} \sim (F_T/f_\pi)\sqrt{3/N_T}\, m_\rho \sim 500 \text{ GeV}\sqrt{\frac{6}{N_T N_D}}$$

The above estimate has been done by scaling QCD and taking $v_{12} = 100$ GeV , $f_\pi = 125$ MeV.

Note that, in order to generate the Yukawa matrix $Y_{12}^{u,d}$, the SM and TC fermions have to be embedded in larger multiplets that transform under the so called extended technicolor (ETC) group. The ETC gauge group breaks into SM $\otimes$ TC group at the ETC scale $M_{\rm ETC}$. Integrating out the gauge bosons corresponding to the broken generators introduces four-fermion operators like,

$$(g^2_{\rm ETC}/M^2_{\rm ETC})\,(\bar{q}q)(\bar{Q}Q) \tag{16}$$

where q and Q are SM and TC quarks respectively, $M_{\rm ETC}$ is the mass of the gauge boson and $g_{\rm ETC}$ is the gauge coupling. This operator generates a mass term for q when, at a lower scale, the TC group becomes strongly coupled and the quark Q forms a condensate,

$$m_q \sim (g^2_{\rm ETC}/M^2_{\rm ETC})\,\langle\overline{Q}Q\rangle_{\rm ETC}. \tag{17}$$

$\langle\overline{Q}Q\rangle_{\rm ETC}$ is the value of the technifermion condensate at the ETC scale which can be related to its value at the TC scale by,

$$\langle\overline{Q}Q\rangle_{\rm ETC} = \langle\overline{Q}Q\rangle_{\rm TC}{\rm Exp}\left[\int_{\Lambda_{\rm TC}}^{M_{\rm ETC}} d\mu/\mu\,\gamma_m(\mu)\right].$$

where $\langle\overline{Q}Q\rangle_{\rm TC} = 4\pi F_T^3$ and γ is the anomalous dimension of the operator. Note that the third generation fermions are also promoted to ETC multiplets, but the ETC breaking pattern is such that, like the first generation fermions, they get small masses from the TC sector, thus generating $Y_{12}^{u,d}$ of the form we have assumed. In this setup, therefore, the lightness of the first two generation fermions is a natural consequence of the fact that the dominant contribution to their masses is from higher dimensional operators generated at a high scale.

A nice feature of this technicolor setup is that the all important μ coupling is generated automatically from the yukawa couplings, $Y_3^{u,d}$, of Φ_3. In order to couple Φ_3 to SM fermions, $\Phi_3\bar{q}q$ in an ETC invariant way, we must also couple it to the condensate $\Phi_3\langle\bar{Q}Q\rangle$, thus generating a μ-term in the potential. The top Yukawa coupling, for instance, can generate,

$$\mu \sim \frac{4\pi v_{12}^2}{N_D^{3/2}} \sim \frac{(150\text{ GeV})^2}{(N_D/3)^{3/2}} \tag{18}$$

for $v_{12} \sim 100$ GeV which gives $m_A \sim 250$ GeV. One can also check that, in this setup, quartic terms in

the scalar potential in Eq. (6) would arise from irrelevant operators above the TC scale and would thus be subdominant compared to the μ-term.

Let us now discuss in some detail how the usual phenomenological problems associated with technicolor theories can be evaded in our case. Possibly the main issues with technicolor models are: (1) EW precision tests and (2) FCNCs. The flavor problem of standard technicolor theories is far less severe in our case, because we do not demand that the condensate accounts for the large top mass whereas the tension with the EWPO is ameliorated because of the presence of a light Higgs in our theory. Let us discuss these one by one.

The tension with flavour physics arises because the extended technicolor interactions are also expected to generate four fermion interactions involving only SM fermions which give rise to FCNC interactions, for example, those contributing to ΔM_K. One can estimate the $\Delta S = 2$ effective Lagrangian to be,

$$(g_{\rm ETC}^2/M_{\rm ETC}^2)\,\theta_{sd}^2(\bar{s}\gamma_\mu P_{L,R}d)(\bar{s}\gamma^\mu P_{L,R}d) \quad (19)$$

where, θ_{sd} is the mixing angle governing the $s \to d$ transition and should be of the order of Cabibbo angle. Using the experimental value of ΔM_K [33, 34] one gets the bound,

$$g_{\rm ETC}\sqrt{{\rm Re}(\theta_{sd}^2)}/(M_{\rm ETC}) \lesssim (700\ {\rm TeV})^{-1}\,, \text{ which gives,}$$

$$m_q \lesssim (M_{\rm ETC}/\Lambda_{\rm TC})^{\bar{\gamma}}\,\frac{4\pi F_T^3}{(700\,{\rm TeV})^2{\rm Re}(\theta_{\rm sd}^2)} \quad (20)$$

where, $\bar{\gamma} = \int_{\Lambda_{\rm TC}}^{M_{\rm ETC}} \frac{d\mu}{\mu}\gamma_m(\mu)\,/\,\log(M_{\rm ETC}/\Lambda_{\rm TC})$ is a parameter which depends on the UV details of the strongly coupled sector. Theories with $\bar{\gamma} \approx 0$ are called running TC theories whereas $\bar{\gamma} \to 1$ is called the walking limit. Taking $\bar{\gamma} = 1$, $F_{\rm TC} \sim 100$ GeV and $\theta_{sd} \sim 0.1$ we get $m_q \lesssim 2.5$ GeV for $M_{\rm ETC}/\Lambda_{\rm TC} \sim 1000$. Hence, obtaining the correct charm quark mass is not a problem in the the walking limit. It is also clear that getting the correct top mass would have been impossible, if we did not have an additional Higgs doublet. Values smaller than $\bar{\gamma} = 1$ can be made compatible with data, if flavour symmetries are imposed in the ETC sector [35] so that smaller values of $M_{\rm ETC}$ are allowed. Walking also helps raise the mass of the pseudogoldstone bosons from chiral symmetry breaking (other than those corresponding to $A, H^\pm$) to the TeV scale [35].

We now discuss the constraints from EWPO on the contribution of the technicolor resonances. The contribution of the resonances to the Peskin-Takeuchi S-parameter is known to be positive and the estimate in [31, 32] gives, $S \sim 0.25\,N_T N_D/6$, which assumes that the technicolor theory is a scaled up version of QCD. The usual tension of technicolor theories with EWPO, however, is much milder in our case because we already have a light Higgs in our spectrum. The electroweak fit assuming a 125 GeV Higgs gives at 95% C.L. [36, 37], $S = 0.05 \pm 0.22$, $T = 0.09 \pm 0.26$. We thus see that $S = 0.25$ is still allowed by the fit at 95% C.L., although a contribution to the T parameter of similar magnitude is also required. This can be achieved in various ways and we refer to [35] and the references therein for more details. In fact, adding an additional contribution $\Delta S\,, \Delta T \sim 0.25$ to the infrared contribution enlarges the allowed parameter space, as can be seen from Fig. 1. Note that the absence of a light Higgs boson shifts the allowed region in the $S-T$ plane by $\Delta S \sim -0.15$ and $\Delta T \sim 0.20$ [38] making the tension with EWPO much stronger. Going back to Eq. (15), we can now see that in the TC limit the reduction in the Couplings of the Higgs bosons to the light fermions can be even more significant than the 2HDM case. For example, from Fig. 1 we get the following estimates for the signal strengths,

$$\left(\mu_f\right)_{\rm 2HDM,\,TC} \gtrsim\ 0.09\,,\ 0.03\,. \quad (21)$$

Note that in the framework presented above the fact that v_{12} and v_3 are of the same order seemingly gives rise to a coincidence problem. In a more complete model, however, it is possible that that the spontaneous symmetry breaking in the SM-like Higgs sector would trigger the dynamical breaking in the Technicolor-like sector which would make this whole setup more compelling.

In summary, in this letter we construct a model where the masses of the fermions of the first two generations are not associated with the Standard Model Higgs mechanism, but with another subdominant sequestered source of electroweak symmetry breaking e.g., a technicolor sector. This structure can naturally explain the smallness of the mixing between the light and heavy generations as well as potentially the lightness of the first two generations. Furthermore, in the technicolor realisation the smallness of the masses is a natural consequence of the fact that they are induced by irrelevant operators. We find that a reduction of the couplings of the observed Higgs boson to the light fermions is a generic prediction of this setup. Our framework can be experimentally tested by direct detection of the pseudogoldstone states $A, H^\pm$ and the TC resonances. We also predict $\mathcal{O}(10\%)$ deviations of Higgs couplings to gauge bosons and

third generation fermions from Standard Model values which can be measured in future precision Higgs precision measurements. A direct measurement of the couplings of the observed Higgs to light fermions will be possible at the ILC and, in particular, with even greater sensitivity at the TLEP [9].

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Note added While this paper was in the final stage of completion, ref [39] appeared on arXiv with a related proposal, however, with a different motivation and a limited overlap.

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